

Covariant Action for a D=11 Five-Brane with the Chiral Field

Paolo Pasti^{1*}, Dmitri Sorokin^{2†} and Mario Tonin^{1 ‡}

¹*Università Degli Studi Di Padova, Dipartimento Di Fisica “Galileo Galilei”
ed INFN, Sezione Di Padova, Via F. Marzolo, 8, 35131 Padova, Italia*

²*National Science Center
Kharkov Institute of Physics and Technology,
Kharkov, 310108, Ukraine*

Abstract

We propose a complete Born–Infeld–like action for a bosonic 5–brane with a worldvolume chiral field in a background of gravitational and antisymmetric gauge fields of $D = 11$ supergravity. When the five–brane couples to a three–rank antisymmetric gauge field, local symmetries of the five–brane require the addition to the action of an appropriate Wess–Zumino term. To preserve general coordinate and Lorentz invariance of the model we introduce a single auxiliary scalar field. The auxiliary field can be eliminated by gauge fixing a corresponding local symmetry at the price of the loss of manifest $d = 6$ worldvolume covariance. The double dimensional reduction of the five–brane model results in the Born–Infeld action with the Wess–Zumino term for a $D = 10$ four–D–brane.

PACS numbers: 11.15-q, 11.17+y

Keywords: P–branes, duality, supergravity.

*e-mail: pasti@pd.infn.it

†e-mail: kfti@rocket.kharkov.ua, sorokin@pd.infn.it

‡e-mail: tonin@pd.infn.it

1 Introduction

New types of supersymmetric p-branes have attracted recently a great deal of attention in the course of studying dualities in string theory. Among them are Dirichlet branes (or D-branes) [1, 2] and an M-theory [3] five-brane [4] carrying in its worldvolume a self-dual (or chiral) two-form gauge field [5]. Complete κ -invariant actions for super-D-branes were constructed very recently in [6], while getting the action for the M-theory five-brane remains a challenging problem. Its solution would allow to shed new light on the structure of D=11 M-theory itself. Equations of motion of a D=11 super-five-brane in a superfield form were obtained in [7] by use of a doubly supersymmetric geometrical approach to describing extended objects [8], while only partial results have been obtained about the structure of bosonic part of the super-five-brane action [9, 10, 11, 12]. The action should be of a higher order in the field strength of the chiral field.

An obstacle to get such an action is caused by the presence of the second-rank antisymmetric gauge field whose field strength is self-dual in the free field limit. When one tries to incorporate the self-duality condition into an action a problem arises with preserving manifest Lorentz invariance of the model. This problem has a rather long history and originates from the electric-magnetic duality and the Dirac monopole problem in Maxwell theory (see [13] and references there in). Rather extensive literature has been devoted to the problem of Lorentz covariance of self-dual field models [14]–[28] in connection with their important role in multidimensional supergravity and string theory. A non-manifestly Lorentz covariant formulation of chiral boson actions has been developed in [14, 15, 16]. As a generalization of this approach a d=6 Born-Infeld-like action for a self-interacting chiral two-form was proposed in [12] as a base for the construction of the M-theory five-brane action. The lack of manifest Lorentz invariance will result in the lack of manifest general coordinate invariance of the worldvolume of the five-brane, which can substantially complicate the construction and study of the complete super-five-brane action in such a formulation. Thus a covariant formulation which would reproduce the results of [12] in a noncovariant gauge is desirable.

The purpose of the present article is to present such a formulation. We propose a $d = 6$ worldvolume covariant action of a Born-Infeld type for a five-brane with a chiral boson in its worldvolume. The action is a generalization of a Lorentz covariant formulation of duality symmetric and self-dual fields proposed in [26, 27]. In contrast to a covariant formulation of chiral bosons with infinitely many auxiliary fields [18, 19, 21, 22, 23, 24] in our approach $d = 6$ covariance is achieved by introducing a single auxiliary scalar field entering the action in a nonpolynomial way [26, 27]⁴. The auxiliary field ensures not only worldvolume covariance but also that all the constraints of the model are of the first class [27] which is important for performing covariant quantization. Upon eliminating the

⁴See ref. [27] on a relation between the two approaches and ref. [29] for the application of the approach with infinitely many fields to studying the duality of Born-Infeld actions.

auxiliary scalar field by a noncovariant gauge fixing of a local symmetry the action reduces (in the case of the flat worldvolume metric) to the non-manifestly Lorentz invariant action of [12] describing a self-interacting second-rank chiral field. Since the bosonic p-branes can live in target space of any dimension higher than $D = p + 1$ we do not specify the dimension of the target space of the five-brane at hand, but imply that $D = 11$ having in mind possible supersymmetric generalization of the model.

We couple the five-brane to antisymmetric gauge fields of $D = 11$ supergravity. A remarkable feature of the model is that when the five-brane couples to a three-rank antisymmetric gauge field, local symmetries of the five-brane require the addition to the action of an appropriate Wess–Zumino term. Upon the double dimensional reduction of the five-brane worldvolume from $d = 6$ to $d = 5$ and target space from $D = 11$ to $D = 10$ the five-brane action reduces to the Born–Infeld action with a Wess–Zumino term for a Dirichlet four-brane.

2 The action

Consider a five-brane described by the following action invariant under $d = 6$ general coordinate transformations:

$$S = \int d^6x \left[\sqrt{-g} \frac{1}{4\partial_r a \partial^r a} \partial_m a(x) F^{*mnl} F_{nlp} \partial^p a(x) + \sqrt{-\det(g_{mn} + i\tilde{F}_{mn})} \right], \quad (1)$$

where x^m ($m=0,1,\dots,5$) are local coordinates of the worldvolume,

$g_{mn}(x) = \partial_m X^M(x) g_{MN} \partial_n X^N(x)$ is a worldvolume metric induced by embedding into curved target space with the metric $g_{MN}(X)$ parametrized by coordinates X^M ($M,N=0,\dots,D-1$); $F_{mnl} = 2(\partial_l A_{mn} + \partial_m A_{nl} + \partial_n A_{lm})$ is the field strength of an antisymmetric worldvolume gauge field $A_{mn}(x)$; $g = \det g_{mn}$; F^{*lmn} is the dual field strength:

$$F^{*lmn} = \frac{1}{6\sqrt{-g}} \varepsilon^{lmnpqr} F_{pqr}$$

and

$$\tilde{F}_{mn} \equiv \frac{1}{\sqrt{\partial a}^2} F_{mnl}^* \partial^l a(x).$$

The scalar field $a(x)$ ensures manifest $d = 6$ covariance of the model and is completely auxiliary, as we shall see below.

Note that in spite of the presence of the imaginary unit inside the determinant in (1) the letter is real. This can be seen by rewriting the determinant as a polynomial in powers of \tilde{F} (as in [12])

$$\det(g_{mn} + i\tilde{F}_{mn}) = g \left(1 + \frac{1}{2} \text{tr} \tilde{F}^2 + \frac{1}{8} (\text{tr} \tilde{F}^2)^2 - \frac{1}{4} \text{tr} \tilde{F}^4 \right). \quad (2)$$

Though the argument of the determinant is 6×6 matrix the polynomial in the r.h.s. of (2) stops at the 4-th power of \tilde{F} since (by construction) \tilde{F}_{mn} is degenerate and has rank 4.

If in (1) we take the flat metric and insert the term $-\frac{1}{2}\tilde{F}_{mn}\tilde{F}^{mn}$ instead of the square root the resulting action

$$\begin{aligned} S &= \int d^6x \frac{1}{4(\partial a)^2} \partial_m a(x) F^{*mnl} (F_{nlp} - F_{nlp}^*) \partial^p a(x) \\ &\equiv \int d^6x \left[\frac{1}{24} F_{lmn} F^{lmn} - \frac{1}{8(\partial a)^2} \partial^m a (F - F^*)_{mnl} (F - F^*)^{nlr} \partial_r a \right]. \end{aligned} \quad (3)$$

will describe a free self-dual field A_{mn} with

$$F_{mnl} - F_{mnl}^* = 0 \quad (4)$$

(see [27] for the details).

The action (1) is invariant under worldvolume diffeomorphisms and the following local transformations:

$$\delta A_{mn} = \partial_{[m} \phi_{n]}(x), \quad (5)$$

(which is the ordinary gauge symmetry of the massless antisymmetric fields),

$$\delta A_{mn} = \frac{1}{2} \partial_{[m} a(x) \varphi_{n]}(x), \quad \delta a(x) = 0 \quad (6)$$

and

$$\begin{aligned} \delta a(x) &= \varphi(x), \\ \delta A_{mn} &= \frac{\varphi(x)}{2(\partial a)^2} (F_{mnp} \partial^p a - \mathcal{V}_{mn}), \end{aligned} \quad (7)$$

where

$$\mathcal{V}^{mn} \equiv -2 \sqrt{-\frac{(\partial a)^2}{g}} \frac{\delta \sqrt{-\det(g_{pq} + i\tilde{F}_{pq})}}{\delta \tilde{F}_{mn}}.$$

(The definition of \mathcal{V}^{mn} is chosen in such a way that in the free limit (3) it coincides with \tilde{F}^{mn} .)

The transformations (6) will allow us to algebraically eliminate part of components of the chiral field A_{mn} . The invariance of the action (1) under (5) and (6) is obvious, and the variation of the action under the transformations (7) is

$$\delta S = \int d^6x \frac{1}{2} \left[\sqrt{-\frac{g}{(\partial a)^2}} (F_{mnp} \partial^p a - \mathcal{V}_{mn}) \delta(\tilde{F}^{mn}) + \frac{\sqrt{-g}}{(\partial a)^2} F^{*mnp} F_{mn}{}^q \partial_{[p} a \partial_{q]} \varphi \right] = 0, \quad (8)$$

where the variation of \tilde{F}^{mn} includes the variation of A_{mn} and $a(x)$ with the parameter $\varphi(x)$ (7). The proof that (8) is zero (up to a total derivative) can be performed along the same lines as in [12], where a non-covariant action for a self-interacting chiral field was demonstrated to possess a modified non-manifest Lorentz invariance.

The equations of motion of A_{mn} , which follow from (1) are

$$\varepsilon^{lmnpqr} \partial_n \frac{\partial_p a}{(\partial a)^2} (F_{qrs} \partial^s a - \mathcal{V}_{qr}) = 0. \quad (9)$$

An appropriate gauge fixing of transformations (6) allows one [26, 27] to reduce the general solution of (9) to the form

$$F_{qrs}\partial^s a - \mathcal{V}_{qr} = 0 \quad (10)$$

which is a generalization of the self-duality condition (4) to the case of the self-interacting field A_{mn} .

As in the free field case [26, 27], the equation of motion of $a(x)$ turns out to be a consequence of (9) and, hence, is not a new field equation. This permits, without losing information on the dynamics of the model, to eliminate $a(x)$ directly from the action by gauge fixing transformations (7) ⁵. However, the price for such a gauge fixing is the loss of worldvolume general coordinate invariance, or Lorentz invariance of the model in the flat limit. For instance, in the case of the flat $d = 6$ metric by putting

$$\partial_m a(x) = \delta_m^5 \quad (11)$$

to be the unit vector along the fifth spatial direction of the worldvolume we reproduce the action and equations of motion of a self-interacting chiral field A_{mn} constructed in [12]. Note that in the gauge (11) we can also completely eliminate the components A_{m5} of the gauge field by use of the algebraic local transformations (6):

$$A_{m5} = 0. \quad (12)$$

Then the modified Lorentz transformations of remaining components $A_{\alpha\beta}$ ($\alpha, \beta = 0, 1, \dots, 4$) of Ref. [12] arise in our approach as a combination (which preserves the gauge fixing condition (11)) of the standard Lorentz transformations with parameters $\Lambda_{mn} = -\Lambda_{nm}$ and the transformation (7) :

$$\delta(\partial_m a(x)) = \Lambda_m^n \partial_n a + x^p \Lambda_p^n \partial_n (\partial_m a) + \partial_m \varphi(x) = \Lambda_m^5 + \partial_m \varphi(x) = 0. \quad (13)$$

From (13) it follows that to preserve (11) the parameter $\phi(x)$ of (7) must be of the form:

$$\phi(x) = -x^m \Lambda_m^5 = -x^\alpha \Lambda_\alpha^5. \quad (14)$$

Substituting (11) and (14) into (7) and combining it with the Lorentz transformation mixing $\alpha = (0, 1, \dots, 4)$ directions with the 5 direction we get for $A_{\alpha\beta}$ (in the gauge $A_{m5} = 0$):

$$\begin{aligned} \delta A_{\alpha\beta} &= (x^\gamma \Lambda_\gamma^5) \partial_5 A_{\alpha\beta} - x^5 (\Lambda_\gamma^5 \partial_\gamma) A_{\alpha\beta} - x^\alpha \Lambda_\alpha^5 (\partial_5 A_{\alpha\beta} - \frac{1}{2} \mathcal{V}_{\alpha\beta}) \\ &= (x^\gamma \Lambda_\gamma^5) \frac{1}{2} \mathcal{V}_{\alpha\beta} - x^5 (\Lambda_\gamma^5 \partial_\gamma) A_{\alpha\beta}, \end{aligned} \quad (15)$$

which are exactly the modified Lorentz transformations of Ref. [12]. They coincide with ordinary Lorentz transformations on the mass shell (10). Thus, though upon imposing the gauge fixing condition (11) the generalized self-duality condition loses the manifestly Lorentz covariant form, it is nevertheless invariant under the Lorentz transformations (15).

⁵Note that the gauge $\partial_m a \partial^m a = 0$ (such as, for instance, $a = \text{const}$) is inadmissible because of the presence of $(\partial a)^2$ in the denominator of Eq. (1). This would lead to a singularity of the action.

3 Coupling to D=11 antisymmetric gauge fields

Consider now the propagation of the five-brane in a background of antisymmetric gauge fields of D=11 supergravity. These are a three-rank field $C_{LMN}^{(3)}(X)$ and its dual six-rank field $C_{LMNPQR}^{(6)}(X)$ [30]. The coupling requires the replacement of the field strength $F_{lmn}(x)$ with

$$H_{lmn} = F_{lmn} - C_{lmn}^{(3)} \quad (16)$$

and adding to the action (1) a Wess–Zumino term [10]. The resulting action becomes

$$S = \int dx^6 \left[\sqrt{-\det(g_{mn} + i\tilde{H}_{mn})} + \sqrt{-g} \frac{1}{4\partial_r a \partial_r a} \partial_m a(x) H^{*mnl} H_{mnp} \partial^p a(x) \right] \\ + \int \left[C^{(6)} + \frac{1}{2} F \wedge C^{(3)} \right], \quad (17)$$

where in (16) and (17) the forms $C^{(6)}$ and $C^{(3)}$ are pullbacks into $d = 6$ worldvolume of the corresponding $D = 11$ forms. The last two terms of (17) form the Wess–Zumino term. The coefficient in front of the last term in (17) is singled out by the requirement that the action (17) remains invariant under the transformations (6) and (7), where in the latter F is replaced with H (16). The transformations of the last term in (17) compensate part of the transformations of the second (H^*H) term while the remaining ones are canceled by the corresponding transformations of the Born–Infeld–like part of the action (like in the absence of the background fields). Thus the Wess–Zumino term is required to preserve local symmetries of the action when the five-brane couples to the antisymmetric fields.

Since $C^{(3)}$ and $C^{(6)}$ are $D = 11$ gauge fields transformed as

$$\delta C^{(3)}(X) = 2d\chi(X), \quad \delta C^{(6)}(X) = -d\chi(X) \wedge C^{(3)}(x), \quad (18)$$

for the action (17) to be invariant under (18) the chiral field A_{mn} must transform as:

$$\delta A(x) = \chi(X(x)).$$

This completes the construction of the action for the bosonic 5-brane propagating in the gravitational and gauge field background of 11-dimensional supergravity.

4 Reduction to a four-D-brane in D=10

Let us perform a double dimensional reduction of the action (17) to a four-brane propagating in a D=10 background. This procedure consists in wrapping x^5 dimension of the five-brane around X^{10} and requiring that all the fields of the reduced model do not depend on x^5 and X^{10} . Below we shall consider a simplified variant of the dimensional reduction by putting to zero dilaton and vector field which arise from components of the

dimensionally reduced metrics and postpone a complete and more systematic analysis to more detailed future paper.

To carry out the dimensional reduction we first put the gauge conditions (11) and (12), split the $d = 6$ indices into $d = 5$ and the 5-th ones and drop the index 5. Thus instead of the original three- and six-rank fields we have:

$$F_{lmn} \Rightarrow (F_{\alpha\beta\gamma}, F_{\alpha\beta}), \quad (19)$$

$$C_{lmn}^{(3)} \Rightarrow (C_{\alpha\beta\gamma}^{(3)}, C_{\alpha\beta}^{(2)}), \quad C^{(6)} \Rightarrow 6C^{(5)}. \quad (20)$$

Note that because of the independence of x^5 and by virtue of (12) the field strength $F_{\alpha\beta}$ in (19) is zero, and only the field $A_{\alpha\beta}$ live in the $d = 5$ worldvolume.

In terms of (19) and (20) the reduced action takes the form

$$\begin{aligned} S = \int dx^5 \left[\sqrt{-\det(g_{\alpha\beta} + iH_{\alpha\beta}^*)} + \frac{\sqrt{-g}}{4} H^{*\alpha\beta} H_{\alpha\beta} \right] \\ + \int dx^5 \varepsilon^{\alpha\beta\gamma\delta\sigma} \left[C_{\alpha\beta\gamma\delta\sigma}^{(5)} + \frac{1}{24} F_{\alpha\beta\gamma} C_{\delta\sigma}^{(2)} \right], \end{aligned} \quad (21)$$

where $H^{*\alpha\beta} = -\frac{1}{6\sqrt{-g}} \varepsilon^{\alpha\beta\gamma\delta\sigma} (F - C^{(3)})_{\gamma\delta\sigma}$.

We see that the second term of (17) became of the Wess–Zumino type in (21) and contributes to the Wess–Zumino term, the resulting action being of the form:

$$\begin{aligned} S = \int dx^5 \sqrt{-\det(g_{\alpha\beta} + iH_{\alpha\beta}^*)} + \\ + \int dx^5 \varepsilon^{\alpha\beta\gamma\delta\sigma} \left[C_{\alpha\beta\gamma\delta\sigma}^{(5)} + \frac{1}{12} F_{\alpha\beta\gamma} C_{\delta\sigma}^{(2)} - \frac{1}{24} C_{\alpha\beta\gamma}^{(3)} C_{\delta\sigma}^{(2)} \right]. \end{aligned} \quad (22)$$

The action (22) is a dual form of the action for a four–D–brane [31]. It can be rewritten in the conventional Born–Infeld form involving the field strength $\hat{F}_{\alpha\beta}$ of a vector field [32] by performing a dualization procedure inverse to that described in [31], so we shall not discuss this point in detail but just note that formally it consists in replacing $F_{\alpha\beta\gamma}$, $C_{\delta\sigma}^{(2)}$ and $C_{\alpha\beta\gamma}^{(3)}$ with their duals:

$$iF_{\alpha\beta}^* \rightarrow \hat{F}_{\alpha\beta} \quad iC_{\alpha\beta}^{(3)*} \rightarrow \hat{C}_{\alpha\beta}^{(2)} \quad C_{\alpha\beta}^{(2)} \rightarrow 12i\hat{C}_{\alpha\beta}^{(3)*},$$

and adjusting coefficients in an appropriate way required by the strict dualization procedure [31]. As a result we get

$$\begin{aligned} S = \int dx^5 \sqrt{-\det(g_{\alpha\beta} + (\hat{F} - \hat{C}^{(2)})_{\alpha\beta})} + \\ + \int dx^5 \varepsilon^{\alpha\beta\gamma\delta\sigma} \left[C_{\alpha\beta\gamma\delta\sigma}^{(5)} + (\hat{F} - \hat{C}^{(2)})_{\alpha\beta} \hat{C}_{\gamma\delta\sigma}^{(3)} - \frac{1}{2} \hat{C}_{\alpha\beta}^{(2)} \hat{C}_{\gamma\delta\sigma}^{(3)} \right]. \end{aligned} \quad (23)$$

The last term of (23) can be included into a redefined $\hat{C}^{(5)}$ and then the Wess–Zumino term in (23) take a canonical formal form $\exp(\hat{F} - \hat{C}^{(2)})(\hat{C}^{(5)} + \hat{C}^{(2)})$ [32].

5 Conclusion

We have constructed the complete Born–Infeld–like action for the bosonic 5–brane carrying the worldvolume chiral field in a background of gravitational and antisymmetric gauge fields of $D = 11$ supergravity. To preserve general coordinate and Lorentz invariance of the model we introduced the single auxiliary scalar field. The auxiliary field can be eliminated by gauge fixing the corresponding local symmetry at the price of the loss of manifest space–time covariance. The double dimensional reduction of the five–brane model results in the Born–Infeld action for the $D = 10$ four–D–brane.

The form of the covariant action (17) and its relation to the D–brane action allows one to hope that it can be generalized to describe embedding of five–brane worldvolume into a superfield background of $D = 11$ supergravity. In other words it should be possible to construct a $d = 6$ covariant κ –symmetric action for a super–five–brane of M–theory analogous to that found for the super–D–branes [6], and which may have the form of (17) with all background fields replaced with corresponding superfields. It would be also of interest to compare the five–brane equations of motion obtained from such an action with covariant superfield equations for the M–theory five–brane proposed in [7], in particular, to find out if there is an analog of the auxiliary scalar field being crucial for the covariance of our model. Other interesting problems are to relate by a dimensional reduction the M–theory five–brane to a type IIA five–brane and a heterotic five–brane in ten dimensions. Work in these directions is in progress.

Acknowledgements. The authors are grateful to I. Bandos, C. Preitschopf and K. Lechner for illuminating discussion. This work was supported by the European Commission TMR programme ERBFMRX–CT96–045 to which P.P. and M.T. are associated. D.S. acknowledges partial support from the INTAS Grants N 93–127 and N 93–493.

Note added. After this paper was sent to hep-th and submitted for publication the authors learned about an article by J. H. Schwarz [33] where the chiral tensor field was coupled to $d = 6$ gravity in the noncovariant formulation. When the antisymmetric field background is switch off and the auxiliary field is gauge fixed by Eq. (11) our model reduces to that of [33].

References

- [1] R. G. Leigh, *Mod. Phys. Lett.* **A4** (1989) 2767.
- [2] J. Polchinski, *Phys. Rev. Lett.* **75** (1995) 4724.
J. Polchinski, Tasi lectures of D–branes, NSF-96-145, hep-th/9611050.
- [3] P. Horava and E. Witten, *Nucl. Phys.* **B460** (1996) 506.
J. H. Schwarz, Lectures on Superstrings and M Theory Dualities, hep-th/9607201.

- [4] M. J. Duff and K. S. Stelle, *Phys. Lett.* **253B** (1991) 113.
R. Güven, *Phys. Lett.* **276B** (1992) 49.
- [5] C. G. Callan, J. A. Harvey and A. Strominger, *Nucl. Phys.* **B367** (1991) 60.
G. W. Gibbons and P. K. Townsend, *Phys. Rev. Lett.* **71** (1993) 3754.
D. Kaplan and J. Michelson, Zero modes for the $D = 11$ membrane and five-brane, hep-th/9510053.
K. Becker and M. Becker, *Nucl. Phys.* **B472** (1996) 221.
- [6] M. Cederwall, A. von Gussich, B. E. W. Nilsson and A. Westerberg, The Dirichlet super-three-brane in type IIB supergravity, hep-th/9610148.
M. Aganagic, C. Popescu and J. H. Schwarz, D-Brane Actions with local kappa symmetry, hep-th/9610249.
M. Cederwall, A. von Gussich, B. E. W. Nilsson, P. Sundell and A. Westerberg, The Dirichlet super-p-branes in ten-dimensional type IIA and IIB supergravity, hep-th/9611159.
E. Bergshoeff and P. K. Townsend, Super D-branes, hep-th/9611173.
M. Aganagic, C. Popescu, and J. H. Schwarz, Gauge-Invariant and Gauge-Fixed D-Brane Actions, hep-th/9612080.
- [7] P. S. Howe and E. Sezgin, Superbranes, hep-th/9607227.
P. S. Howe and E. Sezgin, D=11, p=5, hep-th/9611008.
- [8] I. Bandos, P. Pasti, D. Sorokin, M. Tonin and D. Volkov, *Nucl. Phys.* **B446** (1995) 79.
- [9] P. K. Townsend, *Phys. Lett.* **373B** (1996) 68.
- [10] O. Aharony, String theory dualities from M-theory, hep-th/9604103.
E. Bergshoeff, M. de Roo, Tomas Ortin, The eleven-dimensional five-brane, hep-th/9606118.
- [11] E. Witten, Five-brane effective action in M-theory, hep-th/9610234.
- [12] M. Perry and J. H. Schwarz, Interacting chiral gauge fields in six dimensions and Born-Infeld theory, hep-th/9611065.
- [13] M. Blagojević and P. Senjanović, *Phys. Rep.* **157** (1988) 233.
- [14] R. Floreanini and R. Jackiw, *Phys. Rev. Lett.* **59** (1987) 1873.
- [15] M. Henneaux and C. Teitelboim, in Proc. Quantum mechanics of fundamental systems 2, Santiago, 1987, p. 79; *Phys. Lett.* **B206** (1988) 650.
- [16] J. H. Schwarz and A. Sen, *Nucl. Phys.* **B411** (1994) 35.

- [17] W. Siegel, Nucl. Phys. **B238** (1984) 307.
- [18] B. McClain, Y. S. Wu, F. Yu, Nucl. Phys. **B343** (1990) 689.
- [19] C. Wotzasek, Phys. Rev. Lett. **66** (1991) 129.
- [20] I. Martin and A. Restuccia, Phys. Lett. **B323** (1994) 311.
- [21] F. P. Devecchi and M. Henneaux, Phys. Rev. **D45** (1996) 1606.
- [22] I. Bengtsson and A. Kleppe, On chiral p-forms, Preprint USITP 96-12, Stockholm, September 1996.
- [23] N. Berkovits, Manifest electromagnetic duality in closed superstring theory, hep-th/9607070.
- [24] N. Berkovits, Local actions with electric and magnetic sources, hep-th/9610134.
N. Berkovits, Super-Maxwell actions with manifest duality, hep-th/9610226.
- [25] P. Pasti, D. Sorokin and M. Tonin, Phys. Lett. **B352** (1995) 59.
- [26] P. Pasti, D. Sorokin and M. Tonin, Phys. Rev. **D52** (1995) R4277.
- [27] P. Pasti, D. Sorokin and M. Tonin, On Lorentz Invariant Actions for Chiral P-Forms, hep-th/9611100.
- [28] H. O. Girotti Poincaré Invariance of a Quantized Duality Symmetric Theory, hep-th/9612059.
- [29] I. Bengtsson, Manifest duality in Born-Infeld theory, hep-th/9612174.
- [30] A. Candiello and K. Lechner, Nucl. Phys. **B412** (1994) 479.
- [31] A. A. Tseitin, Nucl. Phys. **B469** (1996) 51.
- [32] M. Douglas, Branes within branes, hep-th/9512077.
M. B. Green, C. M. Hull and P. K. Townsend, D-brane Wess-Zumino actions, T-duality and the cosmological constant, hep-th/9604119.
- [33] J. H. Schwarz, Coupling a Self-Dual Tensor to Gravity in Six Dimensions, hep-th/9701008.